

Active particle in a very thin interfacial droplet

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A single light-driven Janus particle confined in a very thin oil droplet at an air-water interface displays intriguing dynamics. While laser activation induces rapid horizontal motion (1 mm/s–1 cm/s) by thermal Marangoni flow, the particle exhibits unexpected periodic circular motions or intermittent irregular motions. We show that the trajectories are the result of a coupling between the self-propulsion of the particle and the spatiotemporal droplet thickness changes. We propose a simple model where the properties of the active particle trajectories are governed by capillary forces and torques due to the confinement of the particle in the thin droplet.

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Soft boundaries, such as droplets [1–6], vesicles [7–10], flexible membranes [11–14], and related systems [15], can serve as soft confinements that active particles can deform and reshape. These soft confinements are of interest in different contexts, both at the fundamental level [14] and for practical uses such as enhanced oil recovery [16], biofilm formation [16,17], targeted drug delivery [18], and soft robotics [11,13].

These boundaries also change the behavior of active particle assemblies [19]. For example, vesicles can undergo significant deformation and possible propulsion due to the internal reorganization of the confined active components [7–9], offering design principles for microscale soft robots and synthetic cells. Droplet systems are particularly intriguing because the interface can trap or alter particle dynamics [17,20–24], opening pathways for the development of novel active interfacial materials. The existence of interfaces can also induce Marangoni flows driven by concentration [25–30] or temperature gradients [31,32], offering a significantly more efficient and faster propulsion mechanism compared to self-diffusiophoresis.

Most studies to date have focused on the behavior of active particles and their mobility in bulk fluids; a few studies have also addressed this issue at single interfaces. However, the physical mechanisms underlying the interplay between activity and interfacial effects, in particular interfacial deformation as may occur when strong confinement is present, remain poorly understood. Thin films [33–36], thin droplets on substrates [37], and interfaces [38,39] provide simplified yet representative systems to explore these questions. Such geometries could help clarify the roles of interfacial tension, active forces, and strong confinement and deformation in shaping particle behavior.

In this Letter, we study the dynamic behavior of a light-driven Janus particle (JP) strongly confined in a very thin, lenslike oil droplet at an air-water interface. Our main observation is that Marangoni-driven fast JP motions exhibit diverse temporal features depending strongly on the droplet geometry confining the particle. We show that the underlying mechanisms governing the particle trajectories arise from the dynamic coupling between the JP motions and the droplet deformation, clarifying the role of capillary confinement in setting active particle trajectories and motions.

The self-propelled Janus particles used here are made of polystyrene spheres, of radius $a = 5 \mu\text{m}$, with a half-coated metal layer (see details in the Supplemental Material [40]). These particles were suspended in tetradecane oil. The suspension was then placed onto an aqueous subphase (ultrapure water with 6.8 mM NaCl) in a Petri dish and chopped with a needle to create separated droplets with radii $R = 10\text{--}600 \mu\text{m}$, as illustrated in Fig. 1(a). We here focus primarily on the dynamics of a single JP within a droplet to avoid the complexity of multiple JPs such as aggregation in the thin droplet due to the long-range capillary bridge attraction [41]. The JP generally resides at the center of the droplet.

To turn on the activity of the JP, a laser beam was used [3]. This activity can be tuned by the illumination intensity I of the top-hat profile laser beam (wavelength: 532 nm and radius: $108 \mu\text{m}$) [3]. Initially, the center of the illumination area was aligned with the center of the stationary droplet. The dynamics of the JP and the droplet were visualized using an inverted microscope (Axio Observer, Zeiss) equipped with a camera (Orca Flash 4.0 from Hamamatsu or Phantom v640 from Vision Research Inc.).

The static droplet radial thickness profiles $h(r)$ are shown in Fig. 1(b) and were obtained from the interference fringe pattern of the droplet as detailed below. The droplet is extremely thin and has a steep profile at the periphery of the JP as illustrated in the schematics of Figs. 1(a) and 1(c).

When the relationship $|\gamma_o - \gamma_{ow}| < \gamma_w < \gamma_o + \gamma_{ow}$ holds for the interfacial tensions of γ_o (air-oil), γ_w (air-water), and γ_{ow} (oil-water), the particle-free oil droplets show a lenslike shape with small contact angles $\theta_1 \approx 0.022$ rad and

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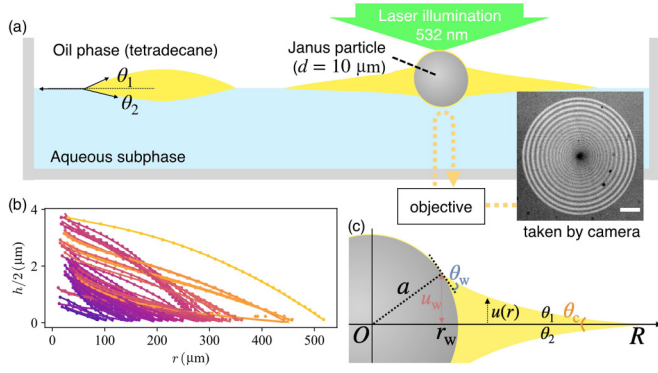


FIG. 1. Experimental setup and thin oil droplet profiles with a Janus particle. (a) Schematics of the interfacial droplets with and without a JP. The inserted image is an example of a droplet with a JP (the black dot in the center). Scale bar = 100 μm . (b) Interfacial profiles $u(r) = h(r)/2$ of droplets with a static JP. The dots correspond to the measured $(r_m, h_m/2)$ values using the fringes, and the solid lines are the fits to Eq. (1). The colors are set by R . (c) Schematic of the lenslike geometry illustrating the contact angle θ_C and apparent contact angle θ_W .

$\theta_2 \approx 0.013$ rad derived from the Young equation for the three fluids, $\gamma_w = \gamma_o \cos \theta_1 + \gamma_{ow} \cos \theta_2$ [Fig. 1(a)]. The particle-free droplets possess a quadratic profile for various radii [Fig. S1(a) of the Supplemental Material [40]]. In contrast, the droplet containing a JP forms a thin droplet with the JP centered, showing a circular symmetry confirmed by the fringes [see inset image of Fig. 1(a)]. The radially averaged thickness profiles $h(r)$ can be reconstructed via fringe counting with the bright ring condition: $h_m = \frac{m\lambda}{2n_o}$ ($m = 1, 2, 3, \dots$ and $\frac{\lambda}{2n_o} = 225$ nm), where λ and n_o are the wavelength of the observation light (around 630 nm) and the refractive index of oil ($n_o = 1.43$), respectively. We assume that the most outward fringe corresponds to the phase difference 2π , and the thickness changes monotonously and continuously. The series of bright spots and corresponding thicknesses (r_m, h_m) give the thickness profile $h(r)$. From the Young-Laplace equation $\gamma \nabla^2 u(r) + p_0 = 0$, we can derive

$$u(r) = -u_0 \ln \frac{r}{R} + u_2 \frac{R^2 - r^2}{2R^2}, \quad (1)$$

which fitted well with all our measured interfacial profiles $(r_m, h_m/2)$. Here, r is the distance from the center of the drop, u_0 is a constant, and $u_2 = p_0 R^2 / \gamma$, where p_0 and γ are the constant internal pressure and the interfacial tension, respectively. Here, we assumed the symmetry of the oil droplet with $\theta_1 = \theta_2$, the profile simplifies to $u(r) = h(r)/2$. As defined in Fig. 1(c), u_0 , u_2 , and R can be determined from the droplet volume V [40], the three-phase contact angles θ_C , and the wetting angle θ_W . The angles θ_C and θ_W are estimated using the profiles in Fig. 1(b) as $\theta_C \simeq 0.58^\circ \pm 0.30^\circ$ and $\theta_W \simeq 23.52^\circ \pm 7.12^\circ$, where the ranges are the standard deviations (see Fig. S1 of the Supplemental Material [40] for the probability distributions).

Since the angles θ_C are material properties determined by the Young relations, the variation may originate from metastability by agitation when creating droplets. Note that θ_W can

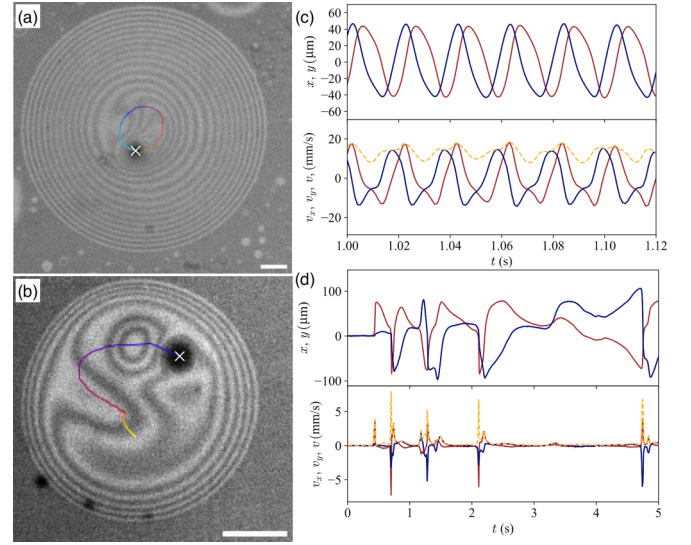


FIG. 2. Examples of regular and irregular motions of JPs. Snapshots of a regular (a) and irregular (b) motion in a droplet under activation corresponding to $t = 1.03$ s and $t = 0.72$ s of panels (c) and (d), respectively, under illumination $I = 426$ W/cm^2 (a) and $I = 415$ W/cm^2 (b). The droplets have radii $R = 239.08$ and 158.13 μm , respectively. The trajectories drawn are for a duration of 20 ms (a) and 0.73 s (b) and are directed from yellow to blue. The JP positions are marked by “ $\times$.” See also Videos S1 and S2 of the Supplemental Material [40]. The scale bars represent 100 μm . (c), (d) Positions $x(t)$, $y(t)$ and velocities $v_x(t)$, $v_y(t)$, in the plane of the interface, corresponding to panels (a) and (b). The red and blue lines are for x and y components, respectively, and the orange dashed lines indicate total speeds $v(t)$.

be understood as the apparent angle because the JP is fully covered by the oil in this study. Therefore, no contact line pinning is expected. The JP has in-plane polarity in a random direction, unlike the interfacial trapped JPs, which have a propensity to have smaller in-plane polarity depending on the surface properties [20,32,42–44].

We now turn to how the thin confining droplets influence the motion of the self-propelled JP. We observe various planar JP fast motions with distinct temporal features using droplets with different sizes under different laser illumination intensities. To comprehend the variety of JP motions, we use the time series of the particle positions $x(t)$, $y(t)$ and velocities $v_x(t)$, $v_y(t)$ in the plane of the interface. Within the droplets, the JP can undergo either regular motions (circular, ellipsoidal, and occasionally back-and-forth motion) with clear periodicity or irregular motions with intermittency (see Fig. 2 and Videos S1 and S2 of the Supplemental Material [40]). In the regular regime, the JP can maintain a periodic motion with a typical velocity of 1 cm/s for more than 17 s (within the limit of our observation time). In contrast, in the irregular regime, the motion of the JP is intermittent, and the trajectory becomes rather chaotic.

To characterize the observed particle trajectories, we examine the power spectral density (PSD) of the velocity $S_{vv}(f) := |\hat{v}_x(f)|^2 + |\hat{v}_y(f)|^2$, where $\hat{v}_j(f) := \frac{1}{N} \sum_{k=0}^{N-1} v_j(t_k) e^{-2\pi i f t_k}$ ($j = x, y$). Here, f is the frequency and t_k is the time, with N being the total number of measurements.

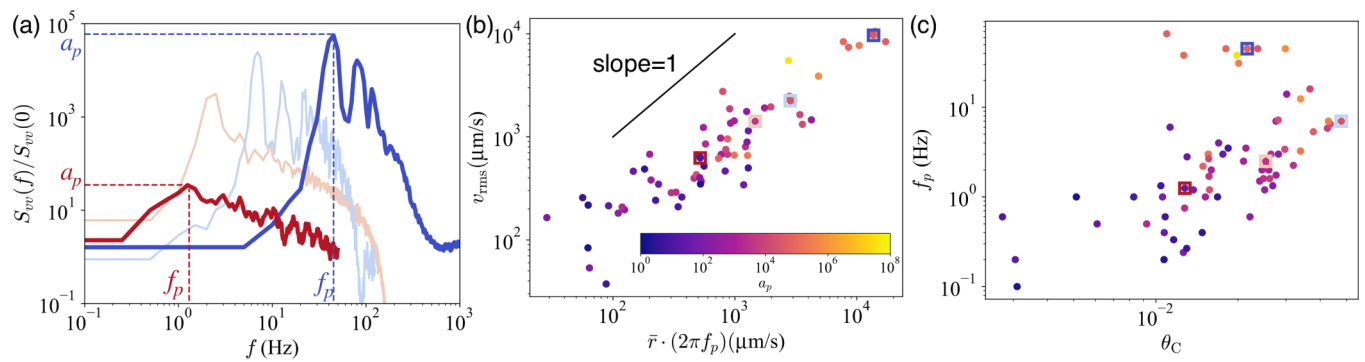


FIG. 3. Analysis of various trajectories in steady states. (a) Typical examples of PSDs. Regular (blue): $I = 426 \text{ W/cm}^2$ corresponding to Figs. 2(a) and 2(c); irregular (red): $I = 415 \text{ W/cm}^2$ corresponding to Figs. 2(b) and 2(d). The other two curves represent intermediate cases in terms of the peak heights ($I = 565$ and 511 W/cm^2). The characteristic frequencies f_p and corresponding amplitudes a_p of the main peaks of the spectra are indicated by dashed lines. (b) Relationship between the rms velocity v_{rms} and $\bar{r} \cdot (2\pi f_p)$. (c) Correlation between the frequency f_p and the contact angle θ_C . In panels (b) and (c), the corresponding cases of panel (a) are marked by the square of the same colors.

Note that we here focus on the steady-state behaviors of JPs within droplets with a roughly constant radius during the entire trajectory. The power spectra of the velocity time series of different trajectories show a prominent peak at a fundamental frequency f_p with a relative amplitude $a_p := S_{vv}(f_p)/S_{vv}(0)$ as shown in Fig. 3(a), satisfying $a_p \sim f_p^2$ (see details in the Supplemental Material [40]). However, two types of spectra can be distinguished: Regular trajectories generally give rise to spectra with well-defined sharp peaks with large amplitudes, while irregular motions give rise to broad spectra with smaller amplitudes. We find a relationship $v_{rms} \approx \bar{r} \cdot (2\pi f_p)$, which is valid for both regular and irregular motions across all our observations.

Here, $v_{rms} := \sqrt{|\mathbf{v}(t)|^2}$ represents the time-averaged speed during steady states, while $\bar{r}$ represents the mean radius of the trajectory. Each JP-droplet combination has different values of I , R , θ_C , the contact angle, and θ_w , the apparent contact angle near the particle, raising the question of which parameters are essential in determining whether the JP motion is regular or irregular. Here, we further find a correlation between the frequency f_p and the contact angle θ_C as shown in Fig. 3(c).

Similarly to light-driven JPs at a planar single interface [31,32], we believe that our JPs are driven by a thermal Marangoni flow at speed $v_p \sim -d\gamma/dT \cdot (\frac{\Delta T}{\eta_o})$, where γ is the interfacial tension, T is the temperature, ΔT is the temperature rise due to light absorption by the metal cap, and η_o is the viscosity of the liquid. This results in velocities of $\mathcal{O}(1 \text{ cm/s})$ for the air-oil interface and $\Delta T \sim 1 \text{ K}$ with the particle moving with its metal cap facing forward as the Marangoni flows go from the hot to the cold region.

To support this, we conducted similar experiments by replacing a JP with a gold core-shell particle [40]. Heating by laser illumination changed the fringe pattern symmetrically around the particle, but resulted in no significant motion. We thus believe that the polarity of the JP is essential, unlike the case of Ref. [45]. In addition, the JP generated flow only locally, as suggested by additional experiments using tracer particles (see Sec. XIV and Video S3 of the Supplemental Material [40]).

Previous studies on thermal Marangoni surfers at an interface showed that the particles propelled at constant speed

$v_p \propto \Delta T/\eta_o \propto I$ at both planar [31,32] and curved interfaces [46,47]. In contrast, the speed of our JPs changes in time (see Fig. 2) through the coupling of the JP motion and the interfacial profile, as seen through the changes of the fringe pattern in Figs. 2(a) and 2(b), and Videos S1 and S2 of the Supplemental Material [40]. Notably, the coupling occurs due to the thin nature of the oil droplets, not due to the intrinsic curvature of the droplet, as demonstrated in our thin film experiments, where $v_p \propto I$ is valid only for a thick film $h \approx 12a$, not in a thin film $h \approx 2a$ (see Sec. VIII of the Supplemental Material [40]).

The coupling ensues as particle motion induces interfacial deformation, and the resulting deformation, in turn, influences the particle dynamics. The physical mechanism of the deformation is due mainly to asymmetric thermal Marangoni flow driving the JP, which hydrodynamically affects the lens deformation due to the finite volume of the oil phase, together with the effect of the moving solid boundary of the JP. This deformation is unique to the nonequilibrium state and can give rise to a local capillary force. In addition, this deformation can yield a net torque on the particle because of the rotational asymmetry of the deformation around the JP, unlike the trivial deformation due to the existence of an off-centered passive particle thicker than the film.

The local capillary force has been proposed in Ref. [48] and is rewritten here in its two-dimensional version:

$$\mathbf{F}(\mathbf{r}) = -\frac{\pi\gamma(h(\mathbf{r}) - h_e)}{1 + (\frac{\partial h}{\partial r})^2} \nabla h, \quad (2)$$

where $\mathbf{r}$ is the position vector and h_e is the equilibrium wetting thickness for the particle.

We utilize Eq. (2) and estimate this capillary force using the reconstructed droplet thickness profile during the steady circular motion. A snapshot of such a profile is shown in Fig. 4(a). The inhomogeneous profile around the JP results in a nonzero net force $\mathbf{F}^{\text{cap}}$ and torque T^{cap} exerted on the JP as defined in the inset of Fig. 4(a). To estimate the force and torque, we assume $h_e := 2u_w$ at $r = r_w$ [see Fig. 1(c)]: the estimated value of wetting thickness without laser illumination. We take a force application line cycle C_ζ with radius ζ centered at the JP. The net force and torque are calculated as $\mathbf{F}_\zeta^{\text{cap}} = \oint_{C_\zeta} \mathbf{F}(s) ds$ and $T_\zeta^{\text{cap}} = \oint_{C_\zeta} \zeta \mathbf{F}(s) \cdot \mathbf{t} ds$, respectively,

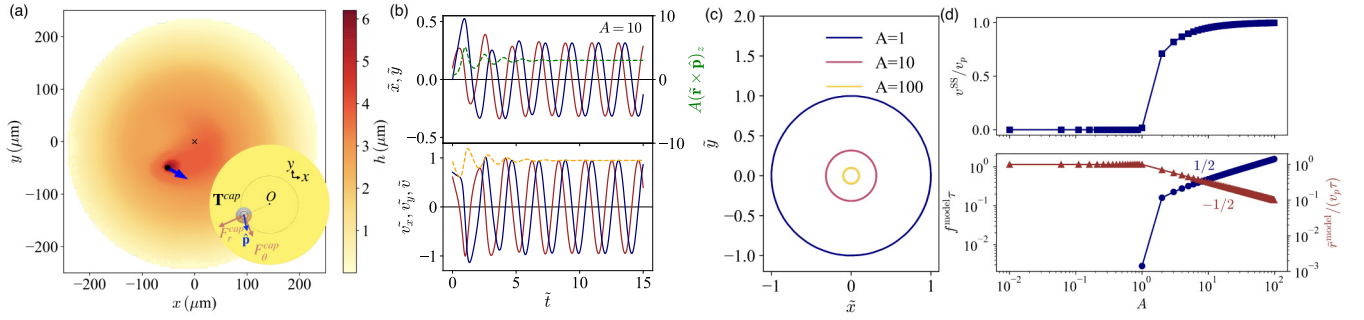


FIG. 4. Capillarity-induced circular motion. (a) Reconstructed thickness profile of a snapshot of the drop with the particle undergoing circular motion ($f_p \approx 45$ Hz, $I = 565$ W/cm²). The black dot and the arrow indicate the particle position and the propelling direction. The center of the droplet is marked as $\times$. Inset: Schematic for the capillary force and torque exerted on the JP in a thin droplet as well as the cap orientation $\hat{\mathbf{p}}$. (b) Examples of the time evolution of position $\tilde{x}(\tilde{t})$, $\tilde{y}(\tilde{t})$ and torque $A\tilde{\mathbf{r}} \cdot \hat{\mathbf{p}}$ (red, blue, and green in the top figure). The corresponding $\tilde{v}_x(\tilde{t})$, $\tilde{v}_y(\tilde{t})$, and speed $\tilde{v}(\tilde{t})$ (orange) are in the bottom figure. For this case, $A = 10$, $\mathbf{r}|_{t=0} = (0.1, 0)$, and $\hat{\mathbf{p}}|_{t=0} = (1/\sqrt{2}, 1/\sqrt{2})$ are used. (c) The limit cycles for different parameters. (d) Steady-state values of the speed v^{SS} , frequency f^{model} , and the radial length $\tilde{r}^{\text{model}}$ for each parameter A (normalized).

where $d\mathbf{s}$ and $\mathbf{t}$ are the infinitesimal length elements along the curve C_ζ and the tangential unit vector along C_ζ , respectively. Since the force field decays with ζ [40], the average within $r_w \leq \zeta \leq \zeta^{\text{max}}$ was taken, where ζ^{max} is a distance above which the forces and torques are negligible. Then, we can obtain the net force $\mathbf{F}^{\text{cap}}$ and net torque T^{cap} . For more details, see Sec. IX and Table S1 of the Supplemental Material [40]. We obtained capillary centripetal forces of the order of 10 nN, which are comparable to the viscous forces $F_v = \xi v_p$ ($\xi = 6\pi a\eta_o$).

Also, there is the capillary positive azimuthal force that pushes the JP forward at the same order of the centripetal force. We often see a thinning at the tail of the JP [as seen also in Figs. 2(a) and 2(b)], so the rear profile becomes steeper, which pushes the JP forward.

Previously, Dauchot *et al.* showed that an active particle subject to a generic self-aligning torque and immersed in a harmonic trap [49] can exhibit different types of trajectories such as orbiting trajectories or so-called climbing trajectories as well as a coexistence between them depending on inertia and the self-alignment intensity. Although the coexistence region occurs only in the presence of sufficient inertia, which is absent in our experiments, some of our observations are in line with this model in particular for the orbiting motion in the capillary confinement. In the scenario of Ref. [49], the particle feels a torque due to the generic self-alignment of the particle. In contrast, the origin of the torque in our case is capillarity and can be estimated experimentally from the dynamic deformation of the drop in the vicinity of the particle. The capillary torque T^{cap} is calculated from the observed profile via Eq. (2) and turns out to be comparable to the viscous torque (see Table S1 in the Supplemental Material [40]). For a particle with counterclockwise circular motion, a positive torque represents an inward rotation of the JP polarity. Note that the nonzero torque T^{cap} means that the capillary force field is not conservative. This is partially because the force [Eq. (2)] is nonlinear. In addition, the equilibrium thickness h_e in Eq. (2) is not necessarily constant and may vary spatiotemporally due to the fluid flow and the asymmetric heating, for example. Therefore, the circular motion can be understood by the combination of the capillary cen-

tripetal force and the capillary positive torque, as well as a self-propulsion.

Here, to better understand the mechanism of the steady circular motion more quantitatively, we consider the following model with an overdamped equation of motion for the particle subjected to rotation by capillarity:

$$0 = -\xi(\dot{\mathbf{r}} - v_p \hat{\mathbf{p}}) - k\mathbf{r}, \quad (3)$$

$$\frac{d\hat{\mathbf{p}}}{dt} = \boldsymbol{\Omega} \times \hat{\mathbf{p}}, \quad (4)$$

$$0 = \mathbf{T}^{\text{cap}} - \xi_r \boldsymbol{\Omega}, \quad (5)$$

where $\mathbf{r}$ is now the position of the particle in the plane, and $\boldsymbol{\Omega}$ and $\hat{\mathbf{p}}$ are the angular velocity of the JP self-rotation and the JP polarity (cap orientation), respectively. Equation (3) takes into account that the particle is powered by the self-propulsion force $F_{\text{sp}} := \xi v_p$, whose value can be shifted by nonzero F_θ^{cap} . The particle also feels the viscous damping and the capillary force in the direction $\mathbf{r}$ or centripetal force, which is approximated, for simplicity and ease, as $-k\mathbf{r}$ with k given by

$$k = 4\pi\gamma \tan^2 \theta_C, \quad (6)$$

derived via the excess interfacial energy by the lowest-order deformation of the droplet profile (see the derivation in the Supplemental Material [40]). This expression matches with Eq. (2) in the limit where the profile is considered locally linear. The centripetal force $|k\mathbf{r}| \sim \mathcal{O}(10 \text{ nN})$ using the measured θ_C is consistent with the estimation from the profile (F_r^{cap}). Experimentally, the capillarity-induced centripetal force can be determined via the relaxation dynamics soon after switching off the laser illumination (see Sec. VII of the Supplemental Material [40]). Neglecting inertia and the azimuthal motion, this force balances the viscous friction, allowing to obtain direct estimates of the value of the spring constant k .

Equations (4) and (5) take into account the capillary torque $\mathbf{T}^{\text{cap}}$ that can rotate the particle and thus its polarity. This torque is balanced by viscous damping with a rotational damping coefficient ξ_r . Here, we may assume that the capillary torque satisfies $\mathbf{T}^{\text{cap}} = \Gamma(\mathbf{r} \times \hat{\mathbf{p}})$ with the coefficient $\Gamma > 0$ for

the observed positive torque. It can be natural to assume the magnitude of the capillary torque to be proportional to that of the capillary force and zero torque for a strictly radial $\hat{\mathbf{p}}$. Equations (3) and (4) are recast as the following nonlinear equations with the nondimensionalized time $\tilde{t} := \tau t$ and position $\tilde{\mathbf{r}} := \mathbf{r}/v_p\tau$ with a single control parameter $A := \frac{v_p\tau^2\Gamma}{\xi_r}$, where $\tau = \xi/k$:

$$\frac{d\hat{\mathbf{p}}}{d\tilde{t}} = A\{(\tilde{\mathbf{r}} \cdot \hat{\mathbf{p}})\hat{\mathbf{p}} - \tilde{\mathbf{r}}\}, \quad (7)$$

$$\frac{d\tilde{\mathbf{r}}}{d\tilde{t}} = \hat{\mathbf{p}} - \tilde{\mathbf{r}}. \quad (8)$$

The nondimensionalized equations are solved using a Python code. The numerical results are shown in Figs. 4(c) and 4(d). Unless the initial polarity is strictly radial, the dynamics results in either steady states with periodic circular motions at a frequency f^{model} with constant torque and speed for $A \geq 1$ or static states without any motion for $A < 1$. In fact, the static states were observed experimentally for a small laser illumination (see Fig. S2 of the Supplemental Material [40]). The steady circular trajectories, namely, the limit cycles, are shown in Fig. 4(c). Large torques, fast propulsion velocities v_p , or large values of τ all result in large values of A , leading to rotation with a smaller radius $\bar{r}$. The steady-state speed $v^{\text{SS}} := \bar{r}^{\text{model}} \cdot (2\pi f^{\text{model}})$ for periodic motions (set to zero for static states) is shown in Fig. 4(d). The speed, comparable to v_p , is consistent with our experimental observations. In addition, the numerical solution leads to the relations $f^{\text{model}} \propto \sqrt{\frac{v_p\Gamma}{\xi_r}}$ and $\bar{r}^{\text{model}} \propto \sqrt{\frac{v_p\xi_r}{\Gamma}}$. Therefore, a larger capillary torque Γ results in motions at a higher frequency f^{model} . Considering the experimental observation that a larger θ_C leads to higher-frequency motion (f_p) as shown in Fig. 3(c), Γ must increase with θ_C . Since $f^{\text{model}} \sim \mathcal{O}(10\text{--}100\text{ Hz})$ for $\tau \sim 10^{-2}\text{ s}$ (for the regular cases), the frequency range turns out to be consistent with our measurements.

This model is consistent with the model proposed by Dauchot *et al.* [49] describing circular motions. However, as we have already pointed out, the speed is not constant but periodic in our circular motions. The thickness profiles at different times during circular motion shown in Figs. S7(a) and S7(b) of the Supplemental Material [40] are similar but not the same; the approximation of the constant capillary force and torque is not perfect. The difference in the azimuthal capillary forces results in the speed modulation in time. We believe that the differences in thickness profiles with time are present because the timescale for oil film relaxation is comparable to that of the particle dynamics.

In our experimental observations, the irregular state occurs at even smaller values of k than those of circular motions. These irregular motions are far beyond the scope of our

model. The approximation that the capillary effects can be modeled as a constant centripetal force and a constant torque most probably breaks down in this case because the thickness profile changes chaotically with the particle motion. The dynamic film deformation must give rise to a time-dependent potential $U(\mathbf{r}, t) = kr^2/2 + U'(\mathbf{r}, t)$, where $U'(t)$ changes in time (see discussion in Sec. VII and Video S4 of the Supplemental Material [40]). Other limitations of the model arise. The nonconservative capillary torque is created by asymmetric Marangoni flow, but fully determining this deformation requires solving the coupled particle dynamics, hydrodynamics, and temperature field, which is beyond the scope of this work. Further, the existence of the solid particle poses challenges to the thin-film hydrodynamic approaches commonly used in related studies [37,50].

We have observed additional intriguing features such as back-and-forth motions. These may be understood as transient states to a circular motion or a result of out-of-plane rotation of the polarity (see discussion in Sec. VIII of the Supplemental Material [40]). Moreover, it is worth noting that, in a transient case with an apparent planar dilation (increase in R), periodic circular motion turned to irregular motion suddenly (see Video S5 of the Supplemental Material [40]). This suggests that the decrease in θ_C due to the dilation and the oil volume conservation may have driven the system into an irregular state as the droplet became much thinner highlighting the role of temporally changing capillary forces and torques in setting the type of particle dynamics.

In conclusion, we observed novel dynamics of active Janus particles confined in thin interfacial droplets. The coupling between droplet thickness and particle motion induces circular motion and other nontrivial trajectories due to the strong confining geometry. We proposed a simple model where capillarity plays a major role in setting the properties of particle trajectories. Insights from studying active particles in thin films could provide a crucial understanding for developing novel active counterparts of particle-based soft matter systems, such as suspension films, capillary suspensions, and Pickering emulsions. The interplay between interfacial properties and activity is expected to yield unique characteristics in activity-controlled functional materials, where manipulating activity offers new avenues for creating smart and responsive systems.

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Data availability. The data supporting this study's findings are available within the article.

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